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Leptogenesis at Low Scale

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Abstract

A typical problem of the leptogenesis scenario is the mismatch between the maximum reheat temperature implied by gravitino overproduction bound and the minimum temperature required to create thermally the lightest right-handed neutrino. We explore the possibility of baryogenesis via leptogenesis in the presence of low scale mass right-handed neutrino. In such a scenario, right-handed neutrinos are created thermally at low reheat temperatures without relying on non-perturbative production mechanisms. We focus on two specific realizations of the scenario, namely the out-of-equilibrium decay of right-handed neutrinos (Fukugita-Yanagida) and the leptogenesis via the LH_u flat direction (Affleck-Dine). We find that in general, the two scenarios are able to produce the required baryon excess for a reasonable amount of CP violation.

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1 Introduction

Recent experimental results gave overwhelming evidence that neutrinos have small but non-vanishing masses [1]. In the standard model (SM), neutrinos are exactly massless, hence the explanation of neutrino experiments requires Physics beyond the SM. Furthermore, neutrino masses appear to be very small with respect to the other fermions ones. If neutrinos are Majorana particles, it is possible to accommodate small neutrinos masses in the SM by introducing the lepton number violating effective operator [2] $O_{\rm eff} = \alpha_{ij} \ell_i^T \tau_2 \vec{\tau} \ell_j H^T \tau_2 \vec{\tau} H/M$, where ℓ_i and H are the lepton and the Higgs doublet respectively. Here M is the scale where "new Physics" is expected to occur, is usually taken as the Planck or the GUT scale. In the former case, the presence of this lepton number violating operator is motivated by the common belief that gravity does not respect any global quantum number [3, 4], or at least this is what happens for example in black holes and wormholes –no hair theorems. In the latter case ($M = M_{\rm GUT}$), the effective operator arises via the see-saw mechanism [5], when integrating-out the heavy right-handed neutrinos (RHNs hereafter).

On the other hand, our Universe appears to be constituted exclusively of baryons. In order not to spoil the Big Bang Nucleosynthesis (BBN) successful predictions of the observed light elements abundances [6], a small baryon excess have to be present. The required value is quantified by the baryon-to-entropy ratio and is given by $Y_B \equiv n_B/s = (7.2 \pm 0.4) \times 10^{-11}$ [7]. To accomplish successfully their task, baryogenesis scenarios [8] have to satisfy three essential conditions [9], namely: *(i.) Baryon number violation*, *(ii.) C and CP violation*, and *(iii.) Departure from thermal equilibrium*. One particularly appealing scenario is the leptogenesis scenario, where lepton number, produced either by the out-of-equilibrium decay of heavy RHN's [10] or by the decay of a scalar condensate carrying non-zero lepton number [11, 12], is reprocessed to a baryon asymmetry via the sphalerons interactions. Given the experimental evidence that lepton number is violated in neutrino oscillation and the fact that proton decay have not been observed yet, the present experimental situation seems to favor this scenario over the other existing baryogenesis scenarios.

A generic problem of thermal leptogenesis scenarios is the mismatch between the maximum reheat temperature implied by gravitino overproduction and the minimum temperature required to thermally create heavy RHNs $T_{\rm RH} \gtrsim 10^{10}$ GeV. To reconcile these two facts, non-thermal creation of RHNs in a low reheat temperature plasma were considered. These mechanisms, however, involve non-perturbative dynamics and are in general sensitive to inflation models. Furthermore, they lead to even more stringent bounds on the reheat temperature, due to the non thermal production of moduli and gravitinos [13, 14].

The aim of this paper is to address this issue in a different perspective. We will consider the situation where the reheat temperature is low (may be as low as the TeV) and we will only consider thermal production of RHNs. This will naturally lead us to consider a class of see-saw models (that we will subsequently call low-scale see-saw models), where RHNs have TeV masses instead of the conventional unification scale.

The paper is organized as follows. In section 2, we give our main motivation for the scenario. In section 3, we study leptogenesis through the out-of-equilibrium decay of low scale RHNs. In section 4, we turn to the Affleck-Dine scenario. Finally, in section 5, we summarize our conclusions.

2 The gravitino problem vs. thermal leptogenesis

As any unwanted relic, gravitinos represents a potential danger for the thermal history of the Universe. Gravitinos are created predominantly via $2 \to 2$ inelastic scatterings of gluons and gluinos quantas. Their relic density and contribution to the energy density are given by [15]

$$Y_{3/2} = 1.1 \times 10^{-10} \left(\frac{T_{\rm RH}}{10^{10}\ {\rm GeV}}\right) \left(\frac{100\ {\rm GeV}}{m_{3/2}}\right)^2 \left(\frac{m_{\tilde g}}{1\ {\rm TeV}}\right)^2 , \qquad (1)$$

$$\Omega_{3/2}\, h^2 = 0.21 \left(\frac{T_{\rm RH}}{10^{10}\ {\rm GeV}}\right) \left(\frac{100\ {\rm GeV}}{m_{3/2}}\right) \left(\frac{m_{\tilde g}}{1\ {\rm TeV}}\right)^2 , \qquad (2)$$

where $m_{\tilde g}$ denotes the gluino mass. The requirement that, if unstable, their late decay do not disrupt the successful BBN predictions, and if stable, their energy density do not overclose the Universe, put tight constraints on their relic abundance. It has been noted that if $m_{3/2}$ >10 TeV or $m_{3/2} <$ keV, then there is no gravitino problem [16, 17]. These requirements can be relaxed if there is a period of inflation and the constraints apply only on post-inflation abundances. From the expression (1), one sees that the gravitino abundance scales linearly with the reheat temperature, therefore the bound on $Y_{3/2}$ translates onto the following bound on the maximum allowed reheat temperature $T_{\rm RH}$ [18]

$$T_{\rm RH} \lesssim (10^6 - 10^9)\ {\rm GeV\ for}\ m_{3/2} = 100\ {\rm GeV} - 1\ {\rm TeV}\,. \qquad (3)$$

There exists however more stringent bounds on $T_{\rm RH}$ from non-thermal production. For generic supersymmetric inflation models, the bound can be as tight as [14] $T_{\rm RH} \lesssim 10^5 (V^{1/4}/10^{15}\ {\rm GeV})$, where $V^{1/4}$ is the height of the inflationary potential.

Let us now see the constraints on the reheat temperature coming from leptogenesis. In the original see-saw model [5] the mass scale of RHNs is typically of $\mathcal{O}(10^{10} - 10^{15})$ GeV. In addition, the bound on the CP parameter [19] for hierarchical RHN's in thermal leptogenesis implies a lower bound on the mass of the lightest RHN $M_{N_1} \gtrsim 10^{10}$ GeV. Consequently, if the thermal leptogenesis scenario is truly *the mechanism* responsible for the the generation of the Baryon Asymmetry of the Universe (BAU), RHNs of this mass have to be produced after inflation. This means a high reheat temperature, at least as high as the mass of the lightest RHN, *i.e.* 10^{10} GeV, potentially conflicting with the gravitino bound discussed above.

A possible way out to get around this problem is to produce RHNs non-thermally, that is during an efficient preheating phase [20]. Non-thermal production, however, can lead in some cases to even more stringent bounds on the reheat temperature. Indeed, for typical hybrid inflation models, the upper bound on the reheat temperature can be as low as 1 TeV [13].

From the above discussion, it is clear that any compelling solution to this problem will, in one way or another, involve low reheat temperatures. After all, we dont know the thermal history of our Universe before BBN. All we know experimentally is that $T_{\rm RH} \geq T_{\rm BBN} \sim$ MeV. In this paper, we will consider a rather exotic solution to this problem, namely the case for leptogenesis when RHNs have a low scale mass. The first benefit of such an approach is that RHNs can be produced thermally with a low reheat temperature $T_{\rm RH} \sim \mathcal{O}(\text{TeV})$, avoiding thus the creation of dangerous relics, like heavy GUT monopoles, and more importantly suppressing the creation of gravitinos. On theoretical grounds, nothing forbids the mass of RHNs to be of $\mathcal{O}(\text{TeV})$. In fact this situation is encountered in many cases (See for *e.g.* [21, 22, 23, 24]). This is also a typical situation that arises in models where the fundamental scale (the GUT scale and/or the quantum gravity scale) is of $\mathcal{O}(\text{TeV})$. In this case, the Yukawa couplings of RHNs have to be much smaller to produce phenomenologically acceptable light neutrino masses. Such a fine-tuning is stable under radiative corrections because that Yukawa couplings are self renormalizable and is protected by supersymmetry. There remains the question of how such suppressed Yukawa couplings can arise in a concrete model. This can be achieved for example by the mean of some R-symmetry that forbids the bare Yukawa coupling between the left and the right-handed neutrinos. As a result the leading Yukawa couplings will be suppressed by powers of a heavy scale [21, 22]. The Yukawa suppression can be obtained upon integrating-out some heavy field as well [25].

3 Thermal Leptogenesis with TeV scale RHNs

We begin by reviewing the basics of the out-of-equilibrium decay leptogenesis scenario. Consider the Minimal Supersymmetric Standard Model extended by three RHNs, one for each generation. The interactions of the RHNs are given by the following superpotential

$$W_N = Y_{ij} L_i H_u N_j + \frac{1}{2} M_i N_i^2 \tag{4}$$

After integrating-out the RHN and electroweak symmetry breaking, the light neutrinos mass matrix will be given by the familiar see-saw formula

$$m_\nu = -Y^T M^{-1} Y \langle H_u \rangle^2 \,. \tag{5}$$

In this scenario, the RHNs must decay out-of-equilibrium. A measure of the departure from thermal equilibrium is given by the parameter K defined as

$$K \equiv \frac{\Gamma_N}{2\,H}\bigg|_{T=M_N}\,, \tag{6}$$

where Γ_N is the decay rate of RHNs and H is the expansion rate of the Universe. The decay is out-of-equilibrium when $K \lesssim 1$. The final baryon asymmetry reprocessed by sphalerons is given by [26]

$$Y_B \equiv \frac{n_B}{s} = \left(\frac{8n_g + 4n_H}{22n_g + 13n_H}\right)\frac{n_L}{s}\,, \tag{7}$$

where n_g and n_H counts the number of fermion generations and Higgses respectively.

The lepton asymmetry produced by the CP-violating out-of-equilibrium decay of the RHNs can be computed using

$$\frac{n_L}{s} = \kappa\,\frac{\varepsilon}{g_*}\,, \tag{8}$$

where g_* is the effective degrees of freedom and κ is the dilution factor, computed by integrating the relevant set of Boltzmann equations [27, 28]. The parameter ε characterizing CP violation in the RHNs decay, can be defined for each RHN separately as [29]

$$\begin{aligned}
\varepsilon_i \;&\equiv\; \frac{\sum_j \Gamma(N_i \to \ell_j h_u) - \sum_j \Gamma(N_i \to \bar{\ell}_j \bar{h}_u)}{\sum_j \Gamma(N_i \to \ell_j h_u) + \sum_j \Gamma(N_i \to \bar{\ell}_j \bar{h}_u)} \\
&=\; -\frac{1}{8\pi}\frac{1}{(YY^\dagger)_{ii}}\sum_{k\neq i}\mathrm{Im}\left[\{(YY^\dagger)_{ik}\}^2\right]\left[F_V\left(\frac{M_k^2}{M_i^2}\right) + F_S\left(\frac{M_k^2}{M_i^2}\right)\right]
\end{aligned} \tag{9}$$

where F_V and F_S are the contributions of the vertex and self-energy respectively. They are given by

$$F_V(x) = \sqrt{x}\ln\left(1 + \frac{1}{x}\right), \quad F_S(x) = \frac{2\sqrt{x}}{x-1} \tag{10}$$

Now, applying the above formulae to TeV mass RHNs, one immediately sees that, due to the smallness of the Yukawa couplings, the decay of RHNs is automatically out-of-equilibrium. In addition to the decay processes, there can be other competing processes that might bring the RHNs to thermal equilibrium, depleting any pre-existing lepton number. These processes have to be out-of-equilibrium too, *i.e.* $\Gamma \simeq \langle n\sigma v\rangle \ll H$. The first such process is the $\Delta L = 2$ scattering $\ell h_u \leftrightarrow \bar{\ell}\bar{h}_u$, via both s and t channel. Other competing processes may involve the t-(s)quark, such as $Nt(\bar{b}) \leftrightarrow \ell b(\bar{t})$. It turns out that due to the Yukawa coupling suppression all these processes are out-of-equilibrium. Finally, it has been noted [30] that the process

$W^{\pm}W^{\pm} \to \ell^{\pm}\ell^{\pm}$, mediated by virtual left-handed neutrinos can lead to stringent constraints on their masses. In our case, it leads to a very mild constraint.

So far for the out-of-equilibrium conditions, now we concentrate on the CP violation parameter ε. As we have seen previously, due to the smallness of the Yukawa couplings, it is very easy to satisfy the out-of-equilibrium condition, however the resulting CP violation parameter ε is too small. This is due to the fact that the decay rate Γ and the CP-parameter ε are both proportional to the same Yukawa couplings combination. From Eqts (9, 10), one sees that the two contributions to the CP parameter ε are sensible to two completely different patters of RHNs masses. While the vertex contribution F_V is enhanced for large hierarchies, the self-energy contribution F_S is so when RHNs are (quasi-)degenerate [1]. In order to enhance the value of ε, one have to exploit the properties of the two functions F_V and F_S. In the next subsection, we will consider the case where RHNs are nearly degenerate [32, 31]. There exist however another possibility, related to the fact that RHNs masses and soft SUSY breaking A-terms are of the same order *i.e.* $\mathcal{O}(\text{TeV})$.

3.1 Leptogenesis with quasi-degenerate TeV scale RHNs

Consider a model where two out of the three RHNs are quasi degenerate, that is N_1, N_2 and N_3 have masses M_1, $M_2 \sim \mathcal{O}(\text{TeV}) \ll M_3$ respectively. The mass splitting $\delta M_{12} \equiv |M_2 - M_1| = \delta \cdot M_0$, where $M_0 \sim$ TeV). Due to their suppressed Yukawa's, RHNs will be long-lived enough to eventually dominate the Universe before decaying. The condition for RHNs dominance can be written $\Gamma_N \ll \Gamma_\varphi$, where Γ_φ is the decay rate of the inflaton [33]. While, RHNs can hardly dominate the energy density of the Universe because of Pauli blocking, this can happen more easily for their scalar partners the RH sneutrinos. Moreover, due to quantum de Sitter quantum fluctuations [34] and for $H_{\text{inf}} \gg M_N \sim \mathcal{O}(\text{TeV})$, RH sneutrinos become coherent over super-horizon scales and can be considered as classical fields with the constant value (vev) $\langle \widetilde{N}^2 \rangle = 3H_{\text{inf}}^4/8\pi^2 M_N^2$. Therefore if the RH sneutrinos scalar potential is just given by the mass term, they are likely to dominate quickly the energy density of the Universe. Given the above discussion, one can compute the lepton asymmetry produced during the decay of N_2 using

$$\frac{n_L}{s} = \frac{3}{4}\frac{T_{N_2}}{M_2}\varepsilon_2 \,, \tag{11}$$

where T_{N_2} is the decay temperature of N_2's computed by equating the energy density of RHNs with the energy density of the Universe when $H \sim \Gamma_2$. Since N_1 and N_2 are quasi-degenerate, we can safely ignore the vertex contribution to the CP parameter

[1] For perturbation theory to hold, the mass splitting $\delta M_{ik} = |M_i - M_k|$ must satisfy $\delta M_{ik} \gg \Gamma$, where Γ is the decay rate of RHNs, otherwise one can no more trust the perturbative calculation based on Eqts (9,10) and one have to rely on a resummation approach [31]. In the limit of exact degeneracy, the CP parameter vanishes.

$(F_S \gg F_V)$. Using Eqts (9) and (10), we can compute the total CP parameter $\varepsilon \simeq \varepsilon_1 + \varepsilon_2$, giving

$$\varepsilon \simeq \frac{1}{8\pi} \sum_{i=1,2} \frac{1}{(YY^\dagger)_{ii}} \operatorname{Im}\left[\{(YY^\dagger)_{12}\}^2\right] \frac{1}{\delta} \tag{12}$$

A rough estimate of the required degeneracy gives $\delta \sim \mathcal{O}(10^{-6} - 10^{-7})$, and perturbativity is clearly satisfied (See footnote 1 on page 5). Such a degeneracy could be ascribed for example to a flavor symmetry, the parameter δ would then characterize its breaking. In the simplest case, the flavor group G_f is taken as a Z_2 and the RHNs have different parity Z_2 assignments, *i.e.* $N_1 \sim$ odd (even) and $N_2 \sim$ even (odd) under Z_2. The flavor symmetry is broken by the vev of the odd field ψ. Restricting to the 12 block, the resulting mass matrix for the RHNs is

$$M_R \sim M_0 \begin{pmatrix} 1 & \delta/2 \\ \delta/2 & 1 \end{pmatrix} \tag{13}$$

with $\delta/2 \equiv \langle\psi\rangle/\Lambda$. The diagonalization of the mass matrix yields two quasi-degenerate RHNs with a mass-splitting δM_0.

3.2 Leptogenesis from soft SUSY breaking A-terms

In the traditional leptogenesis scenario, the contributions of the soft SUSY breaking A-terms to the CP parameter ε are usually neglected. Indeed, SUSY breaking will induce the following A-terms

$$\mathcal{L}_{\text{soft}} = A_{ij} m_{3/2} \widetilde{L}_i \widetilde{N}_j H_u + \text{h.c.}\,. \tag{14}$$

Let us consider the following vertex diagrams, where in the tri-scalar vertex we put the A-term contribution from Eqt (14) instead of the standard SUSY one.

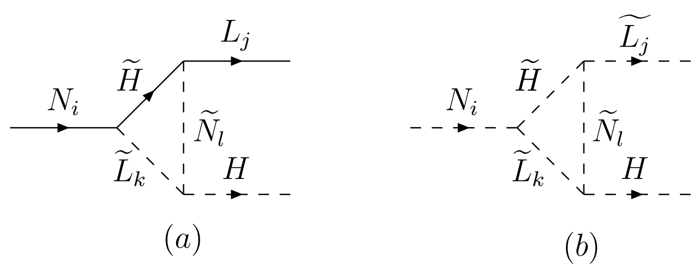

Figure 1: SUSY breaking A-term contributions to the CP parameter ε.

Estimating the contribution of the SUSY soft breaking A-terms to ε and comparing it to the standard SUSY one for each of the two considered diagrams, we obtain

$$
\begin{aligned}
\frac{\varepsilon_{(a)}^{\rm soft}}{\varepsilon^{\rm SUSY}} &\sim |A|\frac{m_{3/2}}{M_i}\sin\delta_{\rm soft}\,, & (15)\\
\frac{\varepsilon_{(b)}^{\rm soft}}{\varepsilon^{\rm SUSY}} &\sim |A|^4\left(\frac{m_{3/2}}{M_i}\right)^4\sin\delta_{\rm soft} & (16)
\end{aligned}
$$

where $\delta_{\rm soft}$ is an effective soft CP phase. From (15), we see that in the conventional leptogenesis scenario, where the mass of the lightest RHN is $M_1 \simeq 10^{10}$ GeV, $\varepsilon_{\rm soft}$ is suppressed with respect to $\varepsilon^{\rm SUSY}$ at least by a factor of 10^{-7}. However, in our scenario, where $M_i \sim m_{3/2}$, the CP asymmetry parameter $\varepsilon_{\rm soft}$ is no more suppressed. It can even dominate the over the SUSY contribution depending on the value of the soft parameters. This means that CP violation may completely originate from the soft SUSY breaking sector, like in the Affleck-Dine case. However, besides enhancing the amount of CP violation, the soft SUSY breaking interactions could bring the RHNs decay at equilibrium, erasing considerably the produced lepton number. A more accurate analysis, requiring the integration of Boltzmann equations, is necessary to reach a firm conclusion.

Finally, it is worth noticing from Eqt. (2) that gravitinos could no more constitute a sizable amount of dark matter in our scenario. Indeed, $\Omega_{3/2}\,h^2 = 0.01-1$, requires the gravitino to be lighter and/or the gluinos masses to be heavier.

4 Affleck-Dine leptogenesis with TeV scale RHNs

Now, we turn to investigate the Affleck-Dine mechanism [11, 12] in the presence of TeV scale RHNs. Consider the LH_u MSSM flat direction given by [2]

$$
L_i = \frac{1}{\sqrt{2}}\begin{pmatrix}\varphi\\ 0\end{pmatrix}, \quad H_u = \frac{1}{\sqrt{2}}\begin{pmatrix}0\\ \varphi\end{pmatrix} \qquad (17)
$$

This flat direction is lifted by the non-renormalizable operator $W_{\rm NR} = \lambda(L\,H_u)^2/M = \lambda\varphi^4/4M$. This operator can be generated via the see-saw mechanism when integrating-out the heavy RHNs. The evolution of the scalar condensate φ in the expanding background is dictated by the classical equation of motion

$$
\ddot{\varphi} + 3H\dot{\varphi} + \frac{\partial V(\varphi)}{\partial\varphi^*} = 0 \qquad (18)
$$

where $V(\varphi)$ is the full potential, including the soft masses, the Hubble induced masses and the A-terms (both from SUSY breaking and the Hubble induced ones)[3].

$$
V(\varphi) = (m_{3/2}^2 - c_H H^2)|\varphi|^2 + a_H H\frac{\varphi^4}{4M}
$$

[2]The factor $\sqrt{2}$ is necessary to have a canonical kinetic term for φ (The Kahler potential is $K = H_u H_u^\dagger + LL^\dagger = \varphi\varphi^\dagger$).

[3]Here, we are simply ignoring thermal effects [35].

$$+a_m m_{3/2}\frac{\varphi^4}{4M} + \text{h.c.} + \frac{|\lambda|^2}{M^2}|\varphi|^6. \tag{19}$$

The constants c_H and a_H depend on the detailed structure of the Kahler potential. In particular the sign of c_H is crucial for the validity of the AD scenario. We assume throughout the paper that it is positive ($c_H > 0$). The evolution of the scalar condensate follows three phases. During inflation, when $H \gg m_{3/2}$ the field φ is over-damped and it settles away from the origin at a distance

$$|\varphi_0| \simeq \left(\frac{c_H M^2 H^2}{|\lambda|^2}\right)^{1/4}. \tag{20}$$

From the last equation, one sees that φ is displaced farther as the neutrino Yukawa coupling λ is smaller. That is why L_i in Eqt. (17) is usually chosen as the neutrino with the smallest Yukawa coupling, L_1 say. When $H \approx m_{3/2}$, the A-terms enter into play and the condensate begins to oscillate. In general, when taking into account thermal effects, the condensate begins to oscillate when the decreasing expansion rate reaches a certain value denoted H_{osc}, determined when the thermal contributions are taken into account [35]. At later times when $H \ll m_{3/2}$, the lepton number is essentially conserved. The evolution of the lepton number, n_L defined as

$$n_L = \frac{i}{2}(\varphi^* \dot{\varphi} - \varphi \dot{\varphi}^*), \tag{21}$$

follows the equation

$$\dot{n_L} + 3Hn_L = \text{Im}\left[\varphi \frac{\partial V(\varphi)}{\partial \varphi}\right] \tag{22}$$

The generated lepton asymmetry can be approximated by integrating the equation (22). This gives

$$n_L \approx \frac{m_{3/2}}{2M}\text{Im}(a_m \varphi^4)t \tag{23}$$

In a matter dominated Universe, the expansion rate scales with time as $H = 2/3t$. Plugging this into the last equation, we get

$$\frac{n_L}{s} \approx \frac{1}{12}\left(\frac{T_{\text{RH}}}{H_{\text{osc}}}\right)\left(\frac{m_{3/2}}{M_*}\right)\left(\frac{M}{M_*}\right)\frac{\delta_{\text{eff}}}{|\lambda|^2}, \tag{24}$$

where $M_* \equiv M_{\text{Planck}}/\sqrt{8\pi} \simeq 2.4 \times 10^{18}$ GeV is the reduced Planck mass and we have dropped constants of $\mathcal{O}(1)$. The effective CP-violating parameter δ_{eff} is defined as

$$\delta_{\text{eff}} \simeq \sin(4 \arg\varphi + \arg a_m) \tag{25}$$

Now, specializing to the low scale see-saw models [22, 21], where Yukawa couplings come-out naturally suppressed as $\lambda \sim |Y^{\text{eff}}|^2 \sim m_{3/2}/M_*$, we get

$$\frac{n_L}{s} \approx \frac{1}{12}\left(\frac{T_{\text{RH}}}{H_{\text{osc}}}\right)\delta_{\text{eff}} \tag{26}$$

Usually the effective CP-violating parameter is assumed to be maximal *i.e.* $\delta_{\rm eff} \simeq 1$. In our case there is no need to do so, since the reheating temperature can be as low as $m_{3/2}$ so TeV mass RHNs are produced thermally, while the gravitinos are not. Typically, the condensate begins to oscillate at $H_{\rm osc} \gtrsim m_{3/2}$. In the extreme case when $T_{\rm RH} \simeq H_{\rm osc} \simeq m_{3/2}$, only a small amount of CP is sufficient to reproduce the observed value, namely $\delta_{\rm eff} \simeq 10^{-9} - 10^{-10}$. Up to now, we did not specify the transmission mechanism of SUSY breaking. We just assumed that some hidden sector will produce the soft breaking scalar masses and A-terms. In the gravity-mediated scenario the A-terms are known to be of the form $a m_{3/2} W + {\rm h.c.}$. This means that $a_m = b_m \lambda$ and $a_H = b_H \lambda$, where b_H, $b_m \sim O(1)$. In this case, the resulting lepton asymmetry is given by

$$\frac{n_L}{s} \approx \frac{1}{12}\left(\frac{T_{\rm RH}}{H_{\rm osc}}\right)\left(\frac{m_{3/2}}{M_*}\right)\delta_{\rm eff} \tag{27}$$

where now the CP violation parameter is defined as

$$\delta_{\rm eff} \simeq \sin\left(4\arg\varphi + \arg a_m + \arg\lambda\right) \tag{28}$$

For the typical values $H_{\rm osc} \sim m_{3/2}$, $T_{\rm RH} \sim 10^9$ GeV and $\delta_{\rm eff} \sim O(1)$, we obtain the right value for the lepton asymmetry.

5 Conclusions

To conclude, motivated by the potential conflict between the gravitino overproduction bound and the high reheat temperature required to produce RHNs thermally, we investigated the baryogenesis through leptogenesis scenario in the presence of low scale RHNs. We have seen that, in such a scenario the Yukawa couplings of RHNs have to be suppressed, in order to give rise to acceptable light neutrino masses. This suppression proved to be useful for many purposes, in particular in satisfying the out-of-equilibrium condition in the FY scenario. Due to this suppression, however, the resulting CP was too small. We used two different mechanisms to enhance the CP parameter: the degeneracy of RHNs and soft A-terms. In the latter case, the necessary CP violation may come entirely from the soft SUSY A-term. We also considered leptogenesis via the LH_u flat direction. We have seen that for generic SUSY breaking scenarios, AD leptogenesis with low scale RHNs is possible, though with reheat temperatures higher than TeV.

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